

Sneutrino-antisneutrino mixing at future colliders

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Abstract. Sneutrino-antisneutrino mixing occurs in a supersymmetric model where neutrinos have nonzero Majorana masses. This can lead to the sneutrino decaying into a final state with a “wrong-sign charged lepton”. In an $e^- \gamma$ collider, the signal of the associated production of an electron-sneutrino and the lighter chargino and their subsequent decays can be $e^- \gamma \rightarrow e^+ \tilde{\tau}_1^- \tilde{\tau}_1^- + \cancel{p_T}$ where the $\tilde{\tau}_1$ s are long-lived and can produce heavily ionizing charged tracks. This signal is free of any Standard Model background, and the supersymmetric backgrounds are small. Such a signal can be experimentally observable under certain conditions which are possible to obtain in an anomaly-mediated supersymmetry breaking scenario. Information on a particular combination of the neutrino masses and mixing angles can also be extracted through the observation of this signal. Sneutrino-antisneutrino mixing at the LHC is currently under study, and asymmetry considerations seem promising there.

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1 Introduction

There has been a tremendous experimental progress in neutrino physics in recent years, and the present data from the solar and atmospheric neutrino experiments contain compelling evidence that neutrinos have tiny masses. It is widely believed that the lepton number (L) may be violated in nature and the neutrinos are Majorana particles. In this case, the smallness of the neutrino masses can be explained by the seesaw mechanism or by dimension-five nonrenormalizable operators with a generic structure. In the context of supersymmetric theories, such $\Delta L = 2$ Majorana neutrino mass terms can induce mixing between the sneutrino and the antisneutrino and a mass splitting ($\Delta m_{\tilde{\nu}}$) between the physical states. The effect of this mass splitting is to induce sneutrino-antisneutrino oscillations, and the lepton number can be tagged in sneutrino decays by the charge of the final state lepton. This can, for example, result in like-sign dilepton signals at e^+e^- colliders and hadron colliders (see the references in [1]). In this talk, based on Ref. [1], we focus on sneutrino-antisneutrino oscillation in the context of an $e^- \gamma$ collider.

In sneutrino-antisneutrino oscillations, which were first discussed in [2], the situation is similar to the flavour oscillation in the $B^0-\bar{B}^0$ system. Suppose the physical sneutrino states are denoted by $|\tilde{\nu}_1\rangle$ and $|\tilde{\nu}_2\rangle$. An initially (at $t = 0$) produced pure $|\tilde{\nu}\rangle$ state is related to the mass eigenstates as

$$|\tilde{\nu}\rangle = \frac{1}{\sqrt{2}}[|\tilde{\nu}_1\rangle + i|\tilde{\nu}_2\rangle]. \quad (1)$$

The state at time t is

$$|\tilde{\nu}(t)\rangle = \frac{1}{\sqrt{2}}[e^{-i(m_1 - i\Gamma_{\tilde{\nu}}/2)t}|\tilde{\nu}_1\rangle + ie^{-i(m_2 - i\Gamma_{\tilde{\nu}}/2)t}|\tilde{\nu}_2\rangle], \quad (2)$$

where the difference between the total decay widths of the two mass eigenstates has been neglected, and the total decay width is set to be equal to $\Gamma_{\tilde{\nu}}$. Since the sneutrinos decay, the probability of finding a “wrong-sign charged lepton” in the decay of a sneutrino should be the time-integrated one and is given by

$$P(\tilde{\nu} \rightarrow \ell^+) = \frac{x_{\tilde{\nu}}^2}{2(1 + x_{\tilde{\nu}}^2)} \times B(\tilde{\nu}^* \rightarrow \ell^+), \quad (3)$$

where the quantity $x_{\tilde{\nu}}$ is defined as

$$x_{\tilde{\nu}} \equiv \frac{\Delta m_{\tilde{\nu}}}{\Gamma_{\tilde{\nu}}}, \quad (4)$$

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and $B(\tilde{\nu}^* \rightarrow \ell^+)$ is the branching fraction for $\tilde{\nu}^* \rightarrow \ell^+$. Here, we assume that sneutrino flavour oscillation is absent and the lepton flavour is conserved in the decay of antisneutrino/sneutrino. If $x_{\tilde{\nu}} \sim 1$ and if the branching ratio of the antisneutrino into the corresponding charged lepton final state is also significant, then one can have a measurable “wrong-sign charged lepton” signal from the single production of a sneutrino in colliders.

It is evident from the above discussion that the probability of the sneutrino-antisneutrino oscillation depends crucially on $\Delta m_{\tilde{\nu}}$ and $\Gamma_{\tilde{\nu}}$. Taking into account the radiative corrections to the Majorana neutrino mass m_{ν} induced by $\Delta m_{\tilde{\nu}}$, one faces the bound [3] $\Delta m_{\tilde{\nu}}/m_{\nu} \lesssim \mathcal{O}(4\pi/\alpha)$. If we consider m_{ν} to be ~ 0.1 eV, then $\Delta m_{\tilde{\nu}} \lesssim 0.1$ keV. Thus, in order to get $x_{\tilde{\nu}} \sim 1$, one also needs the sneutrino decay width $\Gamma_{\tilde{\nu}}$ to be ~ 0.1 keV or so. In other words, this small decay width means that the sneutrino should have enough time to oscillate before it decays. However, such a small decay width is difficult to obtain in most of the scenarios widely discussed in the literature with the lightest neutralino ($\tilde{\chi}_1^0$) being the lightest supersymmetric particle (LSP). In this case, the two-body decay channels $\tilde{\nu} \rightarrow \nu \tilde{\chi}_1^0$ and/or $\tilde{\nu} \rightarrow \ell^- \tilde{\chi}_1^+$ involving the neutralinos ($\tilde{\chi}_1^0$) and the charginos ($\tilde{\chi}_1^+$) will open up. In order to have a decay width $\Gamma_{\tilde{\nu}} \lesssim \mathcal{O}(1)$ keV, these two-body decay modes should be forbidden so that the three-body decay modes $\tilde{\nu} \rightarrow \ell^- \tilde{\tau}_1^+ \nu_{\tau}$ and $\tilde{\nu} \rightarrow \nu \tilde{\tau}_1^{\pm} \tau^{\mp}$ are the available ones. In addition, one should get a reasonable branching fraction for the $\ell^- \tilde{\tau}_1^+ \nu_{\tau}$ final state in order to get the wrong-sign charged lepton signal. It has been pointed out in Ref. [3] that, in order to achieve these requirements, one should have a spectrum

$$m_{\tilde{\tau}_1^{\pm}} < m_{\tilde{\nu}} < m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_1^{\pm}}, \quad (5)$$

where the lighter stau ($\tilde{\tau}_1^{\pm}$) is the LSP. However, having $\tilde{\tau}_1^{\pm}$ as a stable charged particle is strongly disfavoured by astrophysical grounds. This could be avoided, for example, by assuming a very small R -parity-violating coupling which induces the decay $\tilde{\tau}_1 \rightarrow \ell \nu$ but still allows $\tilde{\tau}_1$ to have a large enough decay length to produce a heavily ionizing charged track inside the detector. The spectrum (5) can be obtained in some part of the parameter space in the context of anomaly-mediated supersymmetry breaking (AMSB) [4] with $\Delta m_{\tilde{\nu}} \lesssim \mathcal{O}(4\pi m_{\nu}/\alpha)$. Hence, AMSB seems to have a very good potential to produce signals of sneutrino-antisneutrino oscillation which can be tested in colliders.

The minimal AMSB (mAMSB) model is described by the following parameters: the gravitino mass $m_{3/2}$, the common scalar mass parameter m_0 , the ratio of Higgs vacuum expectation values $\tan \beta$, and $\text{sgn}(\mu)$.

In the mAMSB, assuming that the small neutrino masses are generated by high energy supersymmetry preserving dynamics at an energy scale M far above the weak scale, the relevant part of the superpotential and the soft SUSY breaking interactions are given

by [5]

$$\Delta W_{\text{eff}} = \frac{\Phi_0}{M} \lambda_{ij} (L_i H_2) (L_j H_2), \quad (6)$$

$$\Delta \mathcal{L}_{\text{soft}} = \frac{C_{ij} \lambda_{ij}}{M} (\tilde{\ell}_i h_2) (\tilde{\ell}_j h_2), \quad (7)$$

where H_2 is the Higgs doublet superfield giving masses to the up-type quarks and L_i are the lepton doublet superfields, with the scalar components h_2 and $\tilde{\ell}_i$, respectively. Φ_0 is the Weyl compensator superfield, $C_{ij} \approx M_{\text{aux}}$, and λ is a matrix in flavour space. After the electroweak symmetry breaking, ΔW_{eff} gives a neutrino mass matrix $(m_{\nu})_{ij} = \frac{2}{M} \lambda_{ij} \langle h_2 \rangle^2$. Including the $\Delta L = 2$ contribution from Eq. (7) to the sneutrino mass-squared matrix in the AMSB scenario, one obtains the sneutrino mass splitting [5]

$$(\Delta m_{\tilde{\nu}})_{ij} \approx \frac{M_{\text{aux}}}{m_{\tilde{\nu}}} (m_{\nu})_{ij} = \mathcal{O}(4\pi(m_{\nu})_{ij}/\alpha). \quad (8)$$

At the $e^- \gamma$ collider, we are interested in $(\Delta m_{\tilde{\nu}})_{ee} = \frac{4\pi}{\alpha} (m_{\nu})_{ee}$, since we want to produce an electron-sneutrino. The one-loop contribution to the neutrino mass coming from the sneutrino mass splitting can be significant [3]. Writing this total contribution as $(m_{\nu})_{ee} = (m_{\nu})_{ee}^0 + (m_{\nu})_{ee}^1$, we use the constraint $|(m_{\nu})_{ee}| < 0.2$ eV coming from the searches for the neutrinoless double beta decay. Here, $(m_{\nu})_{ee}^0$ is the tree-level value discussed in Eq. (8) and $(m_{\nu})_{ee}^1$ is the one-loop contribution.

The way to obtain very high energy photon beams is to induce laser back-scattering off an energetic e^{\pm} beam [6]. The use of perfectly polarized electron and photon beams maximizes the signal cross section, although, in reality, it is almost impossible to achieve perfect polarizations. For the laser beam, perfect polarization is relatively easy to obtain, and we shall use $|P_L| = 1$. However, the same is not true for electrons or positrons, and we use $|P_b| = |P_{e^-}| = 0.8$ as a conservative choice. Since we want to produce the sneutrino in this study, the e^- should be left-polarized, i.e. $P_{e^-} = -0.8$. In order to improve the monochromaticity of the outgoing photons, the laser and the e^{\pm} beam should be oppositely polarized, which means $P_L \times P_b < 0$.

2 Signal and backgrounds

We are interested in the production process $e^- \gamma \rightarrow \tilde{\nu}_e \tilde{\chi}_1^-$ and then look at the oscillation of the $\tilde{\nu}_e$ into a $\tilde{\nu}_e^*$. The resulting antisneutrino then decays through the three-body channel $\tilde{\nu}_e^* \rightarrow e^+ \tilde{\tau}_1^- \bar{\nu}_{\tau}$ with a large branching ratio. The chargino $\tilde{\chi}_1^-$ subsequently decays into a $\tilde{\tau}_1^-$ and an antineutrino ($\bar{\nu}_{\tau}$). The neutrinos escape detection and give rise to an imbalance in momentum. The signal is then

$$e^- \gamma \rightarrow \tilde{\nu}_e \tilde{\chi}_1^- \rightarrow e^+ + \tilde{\tau}_1^- + \tilde{\tau}_1^- + \cancel{p_T}, \quad (9)$$

where the two $\tilde{\tau}_1^-$ s are long-lived and can produce heavily ionizing charged tracks inside the detector after traversing a macroscopic distance. The positron serves as the trigger for the event. We assume that the $\tilde{\tau}_1^-$ decays through a tiny R -parity-violating coupling $\lambda_{233} = 5 \times 10^{-9}$ into charged lepton + neutrino pairs so that a substantial number of events do have a reasonably large decay lengths for which the displaced vertex may be visible.

We select the signal events in Eq. (9) according to the following criteria:

1. The transverse momentum of the positron must be large enough: $p_T^{e^+} > 10$ GeV.
2. The transverse momentum of the $\tilde{\tau}_1^-$ s must satisfy $p_T^{\tilde{\tau}_1} > 10$ GeV.
3. The positron and both the staus must be relatively central, i.e. their pseudorapidities must fall in the range $|\eta^{e^+, \tilde{\tau}_1}| < 2.5$.
4. The positron and the staus must be well-separated from each other: i.e. the isolation variable $\Delta R \equiv \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}$ (where η and ϕ denote the separation in rapidity and the azimuthal angle, respectively) should satisfy $\Delta R > 0.4$ for each combination.
5. The missing transverse momentum $p_T^{\cancel{e}} > 10$ GeV.
6. Both the heavily ionizing charged tracks due to the long-lived staus should have a length ≥ 5 cm.

The signal (9) is free of any Standard Model (SM) backgrounds when the λ_{233} coupling is small. However, there are backgrounds from SUSY processes [1] of which $e^- \gamma \rightarrow \tilde{e}_L^- \tilde{\chi}_1^0 \rightarrow e^+ + \tilde{\tau}_1^- + \tilde{\tau}_1^- + p_T$ is the most important one. If the λ_{233} coupling is larger, then the staus decay more rapidly, in which case there are SM backgrounds coming from $e^- \gamma \rightarrow W^- W^- W^+ \nu_e \rightarrow e^+ \ell^- \ell^- + p_T$ where $\ell = \mu, \tau$. However, these backgrounds are quite small [7].

In Fig. 1, we show our results for the total number of positron events for a machine operating at $\sqrt{s_{ee}} = 500$ GeV with 500 fb^{-1} integrated luminosity after imposing the kinematical cuts discussed above, while satisfying $N_e \geq 5\sqrt{N_e + N_B}$, where N_e is the number of signal events and N_B is the number of background events. The region marked by (A) corresponds to a lighter stau mass of less than 86 GeV (see [1] for the references for different experimental constraints). The area below the line X does not satisfy the mass hierarchy of Eq. (5). Thus, the allowed region in the $(m_0 - m_{3/2})$ plane is the one between the area (A) and the line X. The other experimental constraints which we have used are the mass of the lighter chargino ($m_{\tilde{\chi}_1^\pm} > 104$ GeV), the mass of the sneutrino ($m_{\tilde{\nu}} > 94$ GeV) and the mass of the lightest Higgs boson ($m_h > 113$ GeV). In Fig. 2, we show a similar plot in the $(m_0 - m_{3/2})$ plane for a machine operating at $\sqrt{s_{ee}} = 1$ TeV with other inputs remaining the same.

Let us then discuss the change in the number of events when $(m_{\nu})_{ee}^0$ is varied in such a way that it is consistent with the upper limit of 0.2 eV for the total contribution $(m_{\nu})_{ee}$. For this, we choose a machine operating at $\sqrt{s_{ee}} = 500$ GeV. Evidently, larger val-

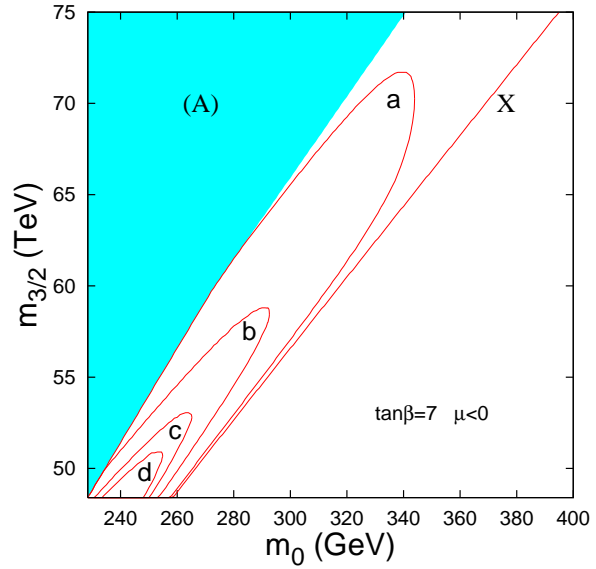


Fig. 1. Parameter regions with $\tan \beta = 7$ and $\mu < 0$. The area (A) represents the parameter region forbidden by the stau mass bound. The mass spectrum (5) is obtained in the region between the area (A) and the line X. Assuming an integrated luminosity of 500 fb^{-1} at $\sqrt{s_{ee}} = 500$ GeV, the numbers of positron events per year inside the contours are: (a) $N_e \geq 50$, (b) $N_e \geq 500$, (c) $N_e \geq 1000$ and (d) $N_e \geq 1300$ for $(m_{\nu})_{ee}^0 = 0.079$ eV so that the total contribution $(m_{\nu})_{ee} \approx 0.2$ eV, while satisfying $N_e \geq 5\sqrt{N_e + N_B}$.

ues of $(m_{\nu})_{ee}^0$ give a larger cross section. This is also shown in Fig. 3 for a sample choice of $m_{3/2} = 50$ TeV, $\tan \beta = 7$ and $\mu < 0$. Assuming an integrated luminosity of 500 fb^{-1} , we have plotted the number of events per year as a function of $m_{\tilde{\nu}_e}$ for different choices of $(m_{\nu})_{ee}^0$. The curves from below correspond to $(m_{\nu})_{ee}^0 = 0.018$ eV, 0.021 eV, 0.035 eV, 0.05 eV, 0.07 eV and 0.081 eV. The corresponding values of the total contribution $(m_{\nu})_{ee}$ are shown in the figure. The horizontal line gives $N_e = 100$ per year. This figure tells us that if we demand the value of N_e to be ≥ 100 , so that the signal significance is $\geq 5\sigma$, then we can probe the value of $(m_{\nu})_{ee}$ down to ≈ 0.05 eV. On the other hand, the current upper limit of 0.2 eV on $(m_{\nu})_{ee}$ sets the upper limit of $(m_{\nu})_{ee}^0 \approx 0.081$ eV. The topmost curve in this figure starts from a slightly higher value of $m_{\tilde{\nu}_e}$, since the bound on $(m_{\nu})_{ee}$ is not satisfied before that. This figure can also be used to extract the value of $(m_{\nu})_{ee}$ with the knowledge of the number of events and other masses.

3 Summary and discussion

We have discussed the potential of an electron-photon collider to investigate the signature of $\tilde{\nu}_e - \tilde{\nu}_e^*$ mixing in an AMSB model which can accommodate $\Delta L = 2$ Majorana neutrino masses. A very interesting feature of such models is that the sneutrino-antisneutrino mass splitting $\Delta m_{\tilde{\nu}}$ is naturally large and is $\mathcal{O}(4\pi m_{\nu}/\alpha)$.

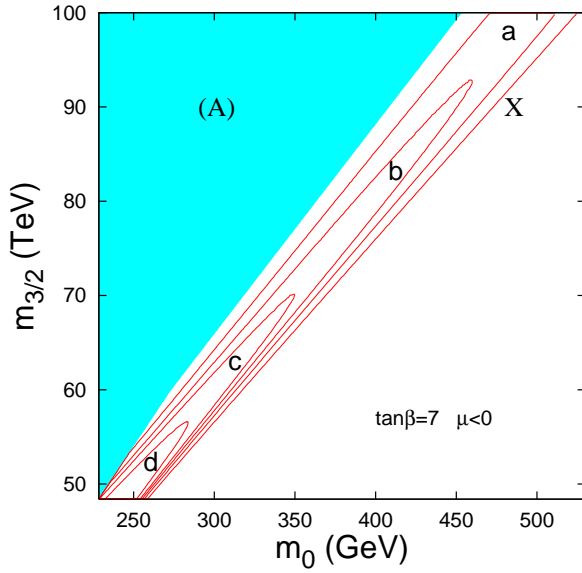


Fig. 2. Parameter regions with $\tan \beta = 7$ and $\mu < 0$. The area (A) represents the parameter region forbidden by the stau mass bound. The mass spectrum (5) is obtained in the region between the area (A) and the line X. Assuming an integrated luminosity of 500 fb^{-1} at $\sqrt{s_{ee}} = 1 \text{ TeV}$, the numbers of positron events per year inside the contours are: (a) $N_e \geq 100$, (b) $N_e \geq 200$, (c) $N_e \geq 300$ and (d) $N_e \geq 500$ for $(m_\nu)_{ee}^0 = 0.079 \text{ eV}$ so that the total contribution $(m_\nu)_{ee} \approx 0.2 \text{ eV}$, while satisfying $N_e \geq 5\sqrt{N_e + N_B}$.

On the other hand, the total decay width of the sneutrino is sufficiently small in a significant region of the allowed parameter space of the model. These two features enhance the possibility of observing sneutrino oscillation signal in various colliders. We have demonstrated that the associated production of the lighter chargino and the sneutrino at an $e^- \gamma$ collider could provide a very clean signature of such a scenario. In addition, this signal can be used to determine $(m_\nu)_{ee}$ which provides important information on a particular combination of the neutrino masses and mixing angles which is not possible to obtain from neutrino oscillation experiments.

Sneutrino-antisneutrino mixing can also be probed in pp collisions at the LHC. The asymmetry between various cross sections can probably indicate whether there is sneutrino oscillation or not. This study is in progress, and the asymmetry considerations seem promising [8].

Acknowledgements

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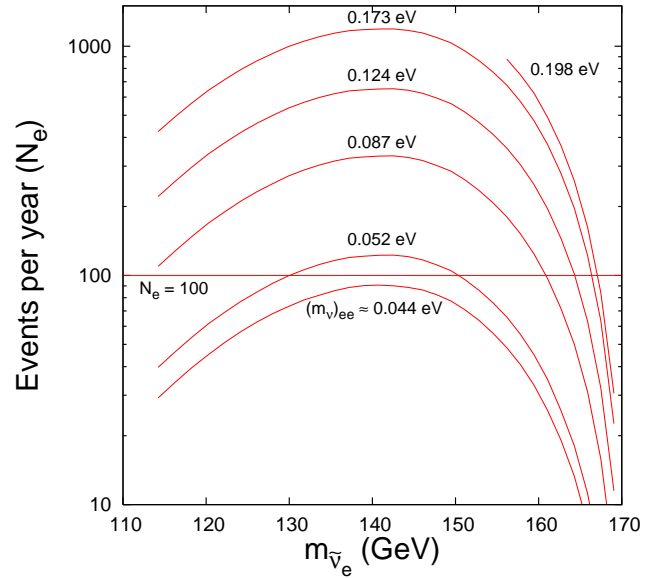


Fig. 3. Number of events (N_e) per year (with an integrated luminosity of 500 fb^{-1}) as a function of $m_{\tilde{\nu}_e}$ for different choices of $(m_\nu)_{ee}^0$ as discussed in the text. Here, $\tan \beta = 7$, $\mu < 0$ and $m_{3/2} = 50 \text{ TeV}$. The values of the total contribution $(m_\nu)_{ee}$ corresponding to each line are shown in the figure. The horizontal line stands for $N_e = 100$ satisfying $N_e \geq 5\sqrt{N_e + N_B}$.

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